On Higher Spin Theory:
Strings, BRST, Dimensional Reductions

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Abstract.   We briefly review some modern developments in higher spin field theory and their links with superstring theory. The analysis is based on various BRST constructions allowing to derive the Lagrangians for massive and massless higher spin fields on flat or constant curvature backgrounds of arbitrary dimensions.

1. Introduction

The paper under consideration is devoted to a brief review of a modern approach to the description of higher spin (HS) fields propagating on flat and constant curvature backgrounds. This approach is based on the BRST technique and has many common features with string field theory. Such a link between HS field theory and (super)string theory is obviously not just formal since the theory of HS fields [1] is conjectured to be an interesting – though not thoroughly investigated – corner of modern string and M–theory.

Being first introduced as a tool for the quantization of gauge theories, the BRST formalism then turned out to be very powerful for the formulation of field theories which correspond to given first quantized systems. A famous example is string field theory [2] that proved to be very useful for the study of string interactions and various superstring vacua containing nonperturbative objects such as D-branes.

After reviewing the triplet structure of HS fields and the way it arises in the tensionless limit of open bosonic string theory [3] – [4] we shall discuss various relevant BRST constructions [5] – [7] and present a general method for constructing HS massive theories in flat and constant curvature backgrounds. Finally, we shortly describe possible further developments such as the supersymmetry of triplet equations and other open problems.

2. Triplets and tensionless limit of string theory

To start with, let us introduce the Fock space spanned by the annihilation and creation operators α_1^M and $\alpha_{-1}^M := (\alpha_1^M)^+$ satisfying $[\alpha_1^M, \alpha_{-1}^N] = g^{MN}$ where g_{MN} is a flat space metric with the mostly plus signature. A completely symmetric tensor field of rank s can be represented in this Fock space by the vector

$$|\varphi\rangle = \frac{1}{s!}\, \varphi_{M_1 M_2 \cdots M_s}(x)\, \alpha_{-1}^{M_1} \alpha_{-1}^{M_2} \cdots \alpha_{-1}^{M_s}|0\rangle. \tag{1}$$

The physical field $|\varphi\rangle$ must satisfy the mass-shell, the transversality, and the tracelessness constraints, i.e it should be annihilated by the operators

$$\tilde{L}_0 = g^{MN}(P_M P_N + i\,\Gamma_{MN}^L P_L) = P^A P_A - i\,\Omega_A{}^{AB} P_B \tag{2}$$

$$\tilde{L}_1 = P_M \alpha_1^M, \quad \tilde{L}_{11} = \frac{1}{2} g_{MN} \alpha_1^M \alpha_1^N, \tag{3}$$

where the operator P_M is the covariant derivative $P_M := -i(\partial_M + \Omega_M{}^{AB}\alpha_{-1\,A}\alpha_{1\,B})$ while $\Omega_M{}^{AB}$ is the spin connection. Let us note that, although in this section we will be dealing with a flat space, the operators (2) – (3) are presented in the generic form since this form will be used in curved backgrounds.

As is well known, in the standard formulation of HS gauge fields [8] the basic object is a doubly traceless completely symmetric tensor field. This can be provided by the off-shell constraint ‡

$$(\tilde{L}_{11})^2|\varphi\rangle = 0 \tag{4}$$

which is invariant under the gauge transformations $\delta|\varphi\rangle = \tilde{L}_{-1}|\Lambda\rangle$ if the parameter of gauge transformations is restricted to be traceless, $\tilde{L}_{11}|\Lambda\rangle = 0$. We defined $\tilde{L}_{-1} := (\tilde{L}_1)^+$. In order to avoid the off-shell constraints on the gauge transformation parameter and basic field, an alternative formulation of higher spin theory was suggested [9] where the field equations are nonlocal (tensors of mixed symmetry were recently discussed in these terms in [10]). One way to obtain the nonlocal equations is via the triplet structure of HS fields [3] though the triplet provides an "unconstrained" description of the HS fields by itself. More specifically the triplet consists of the completely symmetric fields φ (of rank s), C (of rank $s-1$) and D (of rank $s-2$) satisfying

$$\tilde{L}_0|\varphi\rangle = \tilde{L}_{-1}|C\rangle, \quad |C\rangle = \tilde{L}_1|\varphi\rangle - \tilde{L}_{-1}|D\rangle, \quad \tilde{L}_0|D\rangle = \tilde{L}_1|C\rangle, \tag{5}$$

and the gauge transformation rules

$$\delta|\varphi\rangle = \tilde{L}_{-1}|\Lambda\rangle, \quad \delta|C\rangle = \tilde{L}_0|\Lambda\rangle, \quad \delta|D\rangle = \tilde{L}_1|\Lambda\rangle. \tag{6}$$

One can conclude that a triplet, after the suitable gauge fixing, describe the propagation of fields of spins $s, s-2, s-4, \ldots$ In order to describe the irreducible representations,

‡ The rule to make contact with the notation used in the works [3, 4, 9] is to make the replacements: $|\varphi\rangle \to \varphi$, $\tilde{L}_0 \to -\Box$, $\tilde{L}_1 \to -i\partial\cdot$, $\tilde{L}_{-1} \to -i\partial$, $2\tilde{L}_{11} \to {}'$ and $F_0 \to -i\,\partial\!\!\!/$.

an additional field – the compensator $|\alpha\rangle$ – should be introduced. It transforms via the trace of the gauge transformation parameter $\delta|\alpha\rangle = \tilde{L}_{11}|\Lambda\rangle$ and obeys the equation

$$\tilde{L}_{11}|\varphi\rangle - |D\rangle = \tilde{L}_{-1}|\alpha\rangle \tag{7}$$

Such a compensator allows to describe a single irreducible higher spin field.

The triplet structure described above finds its natural origin in string theory [3]. For instance one can take the BRST charge for bosonic open string §, make the redefinition of bc ghosts as $c_m \to \sqrt{2\alpha'}c_m$, $b_m \to \frac{1}{\sqrt{2\alpha'}}b_m$, for $m \neq 0$ and $c_0 \to \alpha'c_0$, $b_0 \to \frac{1}{\alpha'}b_0$, and then go to the $\alpha' \to \infty$ limit. Alternatively, one can take the tensionless limit in the Virasoro constraints prior to the BRST construction [12]. The resulting BRST charge is

$$Q = c_0\tilde{L}_0 + \sum_{m=1}^{\infty}(c_{-m}\tilde{L}_m + c_m\tilde{L}_m^+ + c_{-m}\,c_m\,b_0) \tag{8}$$

where the redefined Virasoro generators satisfy the algebra $[\tilde{L}_m, \tilde{L}_n^+] = \delta_{mn}\tilde{L}_0$. If we restrict the string field to the most general level s leading Regge trajectory

$$|\Phi\rangle = |\varphi\rangle + c_{-1}\,b_{-1}|D\rangle + c_0\,b_{-1}|C\rangle \tag{9}$$

and insert it into the BRST cocycle condition $Q|\Phi\rangle = 0$, then we get the triplet equations (5). The string field gauge transformation $\delta|\Phi\rangle = Q|\Sigma\rangle$ corresponds to the transformation rules (6) if one takes $|\Sigma\rangle = b_{-1}|\Lambda\rangle$. One can consider the general dependence of the string field on an arbitrary finite number of oscillators and recover the generalization of (5), i.e. the generalized triplet [4] which exhausts the bosonic string spectrum in the tensionless limit. Let us also mention that the triplet structure can be successfully deformed to the case of an arbitrary dimensional (A)dS space as well [4].

3. The dimensional reduction

Dimensional reduction might be the most elegant way to produce massive field theories from their known massless counterparts. In this way the massive triplet is recovered, which – besides of having a different field content compared to the usual tensile string equations – is gauge invariant in any dimension rather than only for $D = 26$. The detailed connection of the massive generalized triplet equations and string mass generation mechanism is a still challenging open problem [4]. Here, we concentrate on another side of the dimensional reduction.

It is known that HS fields propagate consistently in flat and constant curvature backgrounds, all of which can be sliced in codimension one constant curvature spaces. Among all corresponding possibilities, three types of dimensional reductions were considered in the literature on higher spins: (i) the usual road going from $\mathbb{R}^{D,1}$ to $\mathbb{R}^{D-1,1}$ [13, 14], (ii) the radial reduction from $\mathbb{R}^{D,1}$ ($\mathbb{R}^{D-1,2}$) to dS_D (AdS_D) recently proposed in [15], and (iii) the reduction from AdS_{D+1} to AdS_D [16]. We will only discuss the first two types of reductions in the BRST framework, nevertheless the procedure we present

§ We follow the notations of [11].

is completely general (in the sense that one could start the reduction from any $D + 1$ dimensional flat or (A)dS background).

The simplest way to make the dimensional reduction and thus generate the masses for triplets is to take the Lagrangian

$$\mathcal{L} = \langle \Phi | \, Q \, | \Phi \rangle, \tag{10}$$

and make x^D coordinate compact, i.e. take the following ansatz for the string field dependence on x^D [14, 17]

$$| \Phi \rangle = V | \Phi' \rangle, \tag{11}$$

where $V = e^{imx^D}$. This is equivalent to the unitary transformation of the BRST charge

$$Q \to V^+ Q V, \tag{12}$$

Obviously the resulting BRST charge is also nilpotent and hermitian in any dimension.

Being the simplest possible one, (12) gives a hint on how to deal with more complicated cases of second class constraints. Indeed, from the point of view of $\mathbb{R}^{D-1,1}$, the operators $L_{\pm 1}$ and $L_0 \equiv P^2 + m^2$ form second class constraints$\|$. A general way to deal with such a problem is therefore to go to a larger phase space where the standard BRST technique applies [6]. In the present illustrative case, we go one dimension up and consider the operators $\tilde{L}_{\pm 1} = P_\mu \, \alpha_{\pm 1}^\mu + P_D \, \alpha_{\pm 1}^D$, $\tilde{L}_0 = P^2 + P_D^2$ which form first class constraints.

Another kind of reduction is needed when the resulting D dimensional space is curved, since in that case the Kaluza-Klein-like reduction discussed above does not eliminate the dependence on the extra coordinate in the D dimensional Lagrangian. Let us consider the procedure of dimensional reduction from flat spaces in its full generality. We initially include the operators $\tilde{L}_{11}$ and $\tilde{L}_{11}^+$ into the set of constraints because they are known to be unavoidable in $(A)dS_D$ space. In the flat $D + 1$ dimensional space with signature $(- + + + ... + \kappa), \kappa = \pm 1$, the constraints – together with the operator $\tilde{G}_0 = g_{MN} \, \alpha_1^{+M} \alpha_1^N + \frac{D+1}{2}$ – form the following algebra

$$[\tilde{L}_1^+, \tilde{L}_1] = -\tilde{L}_0, \quad [\tilde{L}_1^+, \tilde{L}_{11}] = -\tilde{L}_1, \quad [\tilde{L}_{11}^+, \tilde{L}_1] = -\tilde{L}_1^+,$$
$$[\tilde{G}_0, \tilde{L}_1] = -\tilde{L}_1, \quad [\tilde{L}_1^+, \tilde{G}_0] = -\tilde{L}_1^+, \tag{13}$$

with the $SO(2,1)$ subalgebra

$$[\tilde{G}_0, \tilde{L}_{11}] = -2\tilde{L}_{11}, \quad [\tilde{L}_{11}^+, \tilde{G}_0] = -2\tilde{L}_{11}^+, \quad [\tilde{L}_{11}^+, \tilde{L}_{11}] = -\tilde{G}_0. \tag{14}$$

Hermitian conjugation is defined with respect to the integration measure $\sqrt{-\kappa g}$, where $g = det(g_{MN})$.

Next, we construct the nilpotent BRST charge for this system (see [4] – [7] for the details of the construction). The main problem is the presence of the operator $\tilde{G}_0$ which is strictly positive and therefore cannot annihilate the physical state, i.e. one effectively deals with second class constraints. One can deal with the problem in complete analogy

$\|$ Operators with tilde will now correspond to the higher dimensional flat space $\mathbb{R}^{D,1}$, while operators without tilde correspond to $\mathbb{R}^{D-1,1}$ or $(A)dS_D$ space ($M = 0, 1, \ldots, D$ and $\mu = 0, 1, \ldots, D - 1$).

with the case of triplet dimensional reduction. The role of the oscillator α_1^D will be played by an additional oscillator d. First, we build the auxiliary realization of $SO(2,1)$ subalgebra in terms of the oscillator d and a constant parameter h provided the auxiliary realization of $\tilde{G}_0$ depends on this parameter linearly. Second, we define the operators $\hat{G}_0 := \tilde{G}_0 + \tilde{G}_{0\,aux} + h$ and $\hat{L}_{11} := \tilde{L}_{11} + \tilde{L}_{11\,aux}(h)$ as the sum of the initial and auxiliary realizations, and third, we construct the standard BRST charge treating *all* operators as the set of first class constraints

$$Q = c_0 \tilde{L}_0 + c_1 \tilde{L}_1^+ + c_{11} \hat{L}_{11}^+ + c_1^+ \tilde{L}_1 + c_{11}^+ \hat{L}_{11} - c_1^+ c_1 b_0 - b_1^+ c_1^+ c_{11} - c_{11}^+ c_1 b_1$$
$$- c_G(c_1^+ b_1 + b_1^+ c_1 + 2c_{11}^+ b_{11} + 2b_{11}^+ c_{11} + \hat{G}_0 - 3) - c_{11}^+ b_G c_{11}. \tag{15}$$

Finally, we consider an auxiliary phase space (x_h, h) and make the similarity transformation

$$Q \to U^{-1} Q U, \quad U = e^{i\pi x_h} \tag{16}$$

where $\pi = -(\tilde{G}_0 + 2d^+ d - 3 + c_1^+ b_1 + b_1^+ c_1 + 2b_{11}^+ c_{11} + 2c_{11}^+ b_{11})$. Thus, one gets rid of the c_G dependence in the BRST charge while preserving its nilpotency at the same time. Obviously, after the similarity transformation (16) the b_G dependence of Q can be dropped without affecting the nilpotency of the BRST charge.

Now it is possible to eliminate all d^+ dependence in $|\Phi\rangle$ and arrive to triplet and compensator equations (5) and (7) via a partial BRST gauge fixing [4]. If we further gauge away the compensator $|\alpha\rangle$ and eliminate $|C\rangle$ and $|D\rangle$ via their own equations of motion [5], then we obtain the Lagrangian for completely symmetric massless HS fields in arbitrary dimensional flat space $\mathbb{R}^{D,1}$ which generalizes Fronsdal's one [8]

$$\mathcal{L}_{D+1} = \frac{1}{2} \langle \varphi | \, K \, | \varphi \rangle. \tag{17}$$

The kinetic operator K is given by

$$K = -\tilde{L}_0 + \tilde{L}_1^+ \tilde{L}_1 - (\tilde{L}_1^+)^2 \tilde{L}_{11} - \tilde{L}_{11}^+ (\tilde{L}_1)^2 + 2\tilde{L}_0 \tilde{L}_{11}^+ \tilde{L}_{11} + \tilde{L}_{11}^+ \tilde{L}_1^+ \tilde{L}_1 \tilde{L}_{11}. \tag{18}$$

By solving the $D+1$ dimensional double tracelessness constraint (4), we can find the explicit form for the dependence of $|\varphi\rangle$ on the operators a_D^+ [14]. The rank-s field $|\varphi\rangle$ is then given in terms of four D dimensional fields $|\varphi_i\rangle$ $(i = 0, 1, 2, 3)$ of respective ranks $s - i$. The fields $|\varphi_1\rangle$ and $|\varphi_2\rangle$ can be eliminated by fixing the gauge completely.

To address case (i) let us take the dimensional reduction of lagrangian $\mathcal{L}_{D+1}$ to the flat space $\mathbb{R}^{D-1,1}$. To make a connection with the work of Singh and Hagen in four dimensions [18], one should present $|\varphi_0\rangle$ and $|\varphi_3\rangle$ as

$$|\varphi_i\rangle = \sum_{m=0}^{\infty} \frac{(L_{11}^+)^m}{m!} |\varphi_i^{(m)}\rangle, \quad (i = 0, 3), \tag{19}$$

where the fields $|\varphi_i^{(m)}\rangle$ are all traceless, i.e. $L_{11} |\varphi_i^{(m)}\rangle = 0$. This set of fields should correspond to the minimal set of auxiliary fields. The standard procedure to formulate free massive HS field Lagrangians [18] uses the traceless fields of spin s, $s - 2$, $s - 3$, $s - 4$, ... In the approach under consideration, one gets the fields $|\varphi_0^{(0)}\rangle$, $|\varphi_0^{(1)}\rangle$, $|\varphi_0^{(2)}\rangle$, ... of respective ranks s, $s - 2$, $s - 4$, ... and the fields $|\varphi_3^{(0)}\rangle$, $|\varphi_3^{(1)}\rangle$, ... of respective

ranks $s - 3$, $s - 5$, ... which exactly corresponds to [18] at $D = 4$. As a result, we get a possibility to construct a Lagrangian formulation for massive HS field theory in arbitrary dimensional space. The previous procedure should clarify the structure of massive HS field Lagrangian in arbitrary D dimensions.

To address case (ii) let us perform a dimensional reduction of the BRST action (10) where the BRST charge is given by (15). Correspondingly, we now focus on the BRST generalization (including ghosts) of the radial dimensional reduction proposed in [15]. In particular we choose the parameterization of the $D + 1$ dimensional flat space in the form

$$E_M^A = \begin{pmatrix} r\, e_\mu^\alpha & 0 \\ 0 & 1 \end{pmatrix}, \quad E_A^M = \begin{pmatrix} \frac{1}{r}\, e_\alpha^\mu & 0 \\ 0 & 1 \end{pmatrix}. \tag{20}$$

The corresponding tangent space metric is

$$g_{AB} = \begin{pmatrix} \eta_{\alpha\beta} & 0 \\ 0 & \kappa \end{pmatrix}. \tag{21}$$

For the case of the dS_D space embedded into the flat space $\mathbb{R}^{D,1}$ one should take $\kappa = 1$, while for the case of the AdS_D space one should take $\kappa = -1$. The vielbein e_μ^α and its inverse e_α^μ describe in each case the corresponding D dimensional spaces. The measure density of the flat $D + 1$ dimensional space is expressed in terms of the measure density of the D dimensional space as

$$\sqrt{-\kappa\, det(g_{MN})} = r^D \sqrt{-\kappa\, det(g_{\mu\nu})}. \tag{22}$$

The parameterization (20) gives the simple relations between the spin connections:

$$\Omega_\mu{}^\alpha{}_\beta = \omega_\mu{}^\alpha{}_\beta, \quad \Omega_\mu{}^\alpha{}_r = e_\mu^\alpha, \tag{23}$$

the other components of $\Omega_M{}^{AB}$ being zero. Therefore the momentum operator P_M is decomposed as

$$P_\mu = -i(\partial_\mu + \omega_\mu^{\alpha\beta}\alpha_{-1\,\alpha}\alpha_{1\,\beta}) + i\, e_\mu^\alpha(\alpha_{-1\,\alpha}\alpha_1^r - \alpha_{-1}^r\alpha_{1\,\alpha}), \ P_r = -i\,\partial_r. \tag{24}$$

In analogy with the triplet dimensional reduction, let us fix the following r dependence in the state vector by inserting in (12) the following expression for V:

$$V = r^{-(\frac{D+1}{2} + iM + N_0 + N_1)}, \tag{25}$$

where M is a real parameter related to the mass, and the operators $N_0 := c_0 b_0 - b_0 c_0$ and $N_1 := c_1^+ b_1 - b_1^+ c_1$ are ghost number operators. Note that the transformation (12) is no more unitary. However the resulting BRST charge Q is nilpotent and hermitian with respect to the D dimensional integration measure. The important point is that after the transformation (12) the Lagrangian (10) multiplied by the integration measure $dr d^D x \times (22)$ becomes independent from the coordinate $u = \log r$. In other words, the proper dimensional reduction has been done. It could be mentioned that the requirement of hermiticity and independence on u of the resulting D dimensional Lagrangian completely fixes the explicit form of the operator V.

4. Conclusions and perspectives

In this short review we have discussed the connection of massless HS fields with bosonic string spectrum, and the mass generation mechanisms for HS fields via various types of BRST dimensional reductions. Let us note that these results can be extended to the case of the superstring, say type I. For instance, the fermionic triplet could be obtained [4] after the truncation of the string functional to the form

$$|\Phi^R\rangle \;=\; |\Phi_1^R\rangle \;+\; (\gamma_0 \;+\; 2\,c_0\,F_0\,)\,|\Phi_2^R\rangle\,, \tag{26}$$

where γ_0 is bosonic ghost zero mode and $F_0 = P_\mu \psi_0^\mu$. Again, we restrict the functionals $|\Phi_1^R\rangle$ and $|\Phi_2^R\rangle$ to be dependent on only bc ghosts and α_1 oscillators. The mechanism of mass generations could be reproduced following the lines of the bosonic systems. Moreover, supersymmetry of triplets can also be established by making use of the operator W [19] – the fermionic vertex emission operator in CFT language [20] – converting the BRST charge in the R sector to the one in the NS sector

$$Q_R\,W \;=\; W\,Q_{NS}\,, \tag{27}$$

and by avoiding some subtleties – concerning the closure of supersymmetry algebra in NS sector [21] caused by the presence of pictures – thanks to the restriction on the oscillator dependence of the string field (which is possible in the tensionless limit). It seems also interesting to obtain the explicit expressions for massive HS fields in flat and constant curvature backgrounds, and study in detail the spectrum of these models, thus making an appropriate generalization of [18] that seems relevant in the light of modern studies of massive HS fields [15] – [16], [22] – [25]. Another aspect of massive HS field theory is the problem of constructing supersymmetric HS models in flat and (A)dS spaces (the first attempts in this direction have recently been undertaken in [26]).

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